

PREPARED FOR SUBMISSION TO JCAP

LPT-ORSAY-16-17, CPHT-RR013-032016

Re-opening dark matter windows compatible with a diphoton excess

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Abstract. We investigate a simple setup in which an excess in the di-photon invariant mass distribution around 750 GeV, as seen by the ATLAS and CMS collaborations, is originated through a pair of collimated photon pairs. In this framework a scalar state s decays into two light pseudo-Goldstone bosons a , each of which subsequently decays into a pair of collimated photons which are misidentified as a single photon. In a minimal context of spontaneous symmetry breaking, we show that coupling a complex scalar field $\Phi = (s + ia)/\sqrt{2}$ to a fermionic dark matter candidate χ , also responsible for generating its mass, allows for the correct relic density in a large region of the parameter space, while not being excluded by the direct or indirect detection experiments. Moreover, the correct relic abundance can naturally co-exist with a relatively large width for the resonant field s .

Keywords: dark matter theory, particle physics - cosmology connection

ArXiv ePrint: [1603.05601](https://arxiv.org/abs/1603.05601)

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1 Introduction

Recently the ATLAS [1] and CMS [2] collaborations have reported an excess in the di-photon invariant mass distribution in the vicinity of 750 GeV in their 13 TeV data at the level of 10 ± 3 fb [1] and 6 ± 3 fb [2], respectively which can be interpreted as a (pseudo)scalar resonance with a (relatively) large width. The ATLAS has analyzed 3.2 fb^{-1} of data and observed a local significance of 3.9σ for an excess peaked at 750 GeV whereas the CMS, using 3.3 fb^{-1} of data, has reported a local significance of 2.8σ - 2.9σ [3] for an excess peaked at¹ 760 GeV. The global significance is reduced to 2.0σ [4] and $< 1\sigma$ [3] for the ATLAS and CMS, respectively. These announcements have initiated a large amount of works on the subject [5–160], trying to fit the data in effective frameworks [161–168], featuring loop induced (typically by new vector-like fermions, with respect to the standard model (SM)) couplings of the resonance, say s for example, with gluons and photons which are described by dimension five operators of the form:

$$-\mathcal{L}_0 = \frac{c_{BB}}{\Lambda} s B^{\mu\nu} B_{\mu\nu} + \frac{c_{WW}}{\Lambda} s W_i^{\mu\nu} W_{\mu\nu}^i + \frac{c_{GG}}{\Lambda} s G_\alpha^{\mu\nu} G_{\mu\nu}^\alpha, \quad (1.1)$$

where $B_{\mu\nu}$, $W_{\mu\nu}^i$ and $G_{\mu\nu}^\alpha$ are the $U(1)_Y$, $SU(2)_L$ and $SU(3)_C$ field strength, respectively. The parameter Λ represents the scale at which new dynamics appears, which is typically of the order of the mass of new fermions entering in the loop to generate an effective Lagrangian as shown in eq. (1.1). However, it was quickly realized that, in the presence of couplings only with gluons and photons, in order to produce the observed production cross-section $\sim \mathcal{O}(10 \text{ fb})$ and simultaneously a rather large decay width of the resonance, Γ_s , such that $\Gamma_s/m_s \sim 1 - 10\%$ [3, 4] with m_s giving the mass of the resonance, large values of the coupling c_{GG} would be essential. This possibility, however, contradicts (see ref. [169]) with the current

¹ Combining with 19.7 fb^{-1} of data at 8 TeV, the largest CMS excess, assuming a narrow width, i.e., $\sim \mathcal{O}(0.1 \text{ GeV})$, appears at 750 GeV with a local and global significance of 3.4σ and 1.6σ , respectively [3].

observations from di-jet searches $gg \rightarrow s \rightarrow gg$ at the LHC [170, 171], predicting $c_{GG} \lesssim 0.1$ for $\Lambda \sim 2$ TeV. Similarly, high values, $\gtrsim 0.1$, of the coupling c_{BB} are disfavoured², given the constraint from the production through photon fusion [172].

A possible solution to increase Γ_s can be achieved by allowing new decay channels for the resonance. An intriguing possibility is represented by the case in which the additional width is provided by the decay channel into dark matter pairs. Along this line, the authors of refs. [47, 173–175] (see also ref. [176]) have considered a coupling of the resonance with a fermionic (Dirac or Majorana) dark matter candidate as:

$$-\mathcal{L}_1 = g_\chi s \bar{\chi}\chi + m_\chi \bar{\chi}\chi, \quad (1.2)$$

with g_χ as the relevant Yukawa coupling and m_χ representing the mass of the dark matter χ .

However, all these studies converged to show that in the said scenario, the direct detection (in the case of a scalar resonance) or the indirect detection (in the case of a pseudoscalar resonance), strongly constraints the dark matter mass up to $\sim \mathcal{O}(1)$ TeV, due to the concerned values of c_{GG} and c_{BB} couplings, respectively. These associated values, even after respecting the di-jet [169] and photon fusion [172] constraints, appear naturally to accommodate the observed di-photon production cross-section. Combining these limits further with the constraints from monojet searches [177, 178], $gg \rightarrow \bar{\chi}\chi g$, the light dark matter region, $m_\chi \lesssim 375$ GeV, i.e., $m_\chi \lesssim m_s/2$ appears excluded [175].

An interesting alternative has been proposed in refs. [179–182], considering the possibility that the di-photon production process $pp \rightarrow s \rightarrow \gamma\gamma$ can be mimicked by a process of the form $pp \rightarrow s \rightarrow aa \rightarrow 4\gamma$, where a is a light pseudoscalar decaying into two collimated photons. In this scenario a sizable contribution to Γ_s could come from the $s \rightarrow aa$ decay process. This possibility has been investigated, for instance, in ref. [182], where the saa coupling is generated by the spontaneous breaking of a Peccei-Quinn (PQ) symmetry. However, a sizable saa coupling to secure $\Gamma_s = 40$ GeV through $s \rightarrow aa$ decay process is rather unnatural since it would require a small PQ scale $\sim \mathcal{O}(300)$ GeV. Thus, the presence of an invisible branching fraction (Br) has nevertheless been invoked, although the new final states have not been identified as dark matter candidates. The similar and extensive studies have been performed in the case of the next-to-minimal supersymmetric standard model (NMSSM) [183–185]. In this model the LHC signal can be reproduced only for masses $\sim \mathcal{O}(100)$ MeV of the lightest CP-odd Higgs³ since otherwise its branching ratio into two photons would be heavily suppressed compared to the same into SM fermion pairs.

In this work, we will combine the two preceding approaches in the context of a simple model featuring spontaneously broken symmetry. The coupling saa is induced after the spontaneous symmetry breaking by a classical potential term $|\Phi|^4 \rightarrow |(s + ia)|^4$. Thus, the derived saa coupling remains proportional to the quartic coupling. The same complex scalar field Φ is also coupled to the dark matter χ and, is responsible for generating m_χ after the spontaneous symmetry breaking.

We will show that when Γ_s is primarily determined by $s \rightarrow aa$ process, i.e., via $\Gamma(s \rightarrow aa)$, the di-photon excess can be obtained for values of c_{GG} well below 0.1, the limit predicted from the di-jet constraint. Such low values of c_{GG} couplings, as an outgrowth, strongly relax the constraints from monojet and direct dark matter searches while still remain compatible with the WMAP/PLANCK [188, 189] favoured values of the dark matter relic density.

²This scenario corresponds to $c_{WW} = 0$ in eq. (1.1).

³Collimated photons from a light CP-odd Higgs in the NMSSM were previously studied in refs. [186, 187].

Regarding the WMAP/PLANCK constraints, three regions of the parameters space are now open: (i) the electroweak-scale dark matter ($m_\chi \simeq 100$ GeV), (ii) a heavier window ($\gtrsim 400$ GeV, i.e., $\gtrsim m_s/2$) and (iii) a very light case (keV dark matter). For all of these windows, we have checked that the direct and indirect detection limits are respected and are compatible with a relatively large width of the scalar, exchanged as a mediator. Our results are summarized in the final plot where we have shown points respecting all the relevant LHC constraints (e.g., di-jet, monojet, etc.), observations (di-photon production cross-section and large Γ_s) and giving, at the same time, the correct relic abundance for the dark matter.

The paper is organized as follows. After a summary of the model in section 2, we will briefly review under which conditions the process $pp \rightarrow s \rightarrow aa \rightarrow 4\gamma$ can mimic the observed di-photon signal. We will then determine the parameter space allowed by the LHC observations, and show that the di-photon signal can remain compatible with a large width of s in section 3. In the next section we introduce the dark matter candidate and compute its relic abundance, evaluating the parameter space respecting at the same time the observed di-photon excess and the WMAP/PLANCK constraints. Finally we conclude.

2 The Model

The Lagrangian of eq. (1.1) can be expressed for a complex scalar Φ in a similar way. Adding a SM singlet fermion (Majorana or Dirac like) further with eq. (1.1) is straightforward. The complete Lagrangian can then be written as:

$$\mathcal{L} = \mathcal{L}_0 + \mathcal{L}_\Phi + \mathcal{L}_\chi, \quad (2.1)$$

with

$$\begin{aligned} -\mathcal{L}_0 &= \frac{c_{BB}}{\Lambda} \Phi B^{\mu\nu} B_{\mu\nu} + \frac{c_{WW}}{\Lambda} \Phi W_i^{\mu\nu} W_{\mu\nu}^i + \frac{c_{GG}}{\Lambda} \Phi G_\alpha^{\mu\nu} G_{\mu\nu}^\alpha \\ &\quad - \frac{ic_{BB}}{\Lambda} \Phi B^{\mu\nu} \tilde{B}_{\mu\nu} - \frac{ic_{WW}}{\Lambda} \Phi W_i^{\mu\nu} \tilde{W}_{\mu\nu}^i - \frac{ic_{GG}}{\Lambda} \Phi G_\alpha^{\mu\nu} \tilde{G}_{\mu\nu}^\alpha + \text{h.c.}, \end{aligned} \quad (2.2)$$

where $\tilde{B}_{\mu\nu}$, $\tilde{W}_{\mu\nu}^i$, $\tilde{G}_{\mu\nu}^\alpha$ denote the dual field strength, e.g. $\tilde{B}_{\mu\nu} = \frac{1}{2}\epsilon_{\mu\nu\rho\sigma}B^{\rho\sigma}$, and $\Phi = (s + ia)/\sqrt{2}$ is a complex scalar field whose Lagrangian can be expressed as:

$$\mathcal{L}_\Phi = \partial_\mu \Phi \partial^\mu \Phi^* + \mu_\Phi^2 |\Phi|^2 - \lambda_\Phi |\Phi|^4 + \frac{\epsilon_\Phi^2}{2} (\Phi^2 + \text{h.c.}). \quad (2.3)$$

We can see that eq. (2.3) contains a term that explicitly breaks the $U(1)$ symmetry, giving mass to the Goldstone boson a , which is then promoted to a pseudo-Goldstone state. A dark matter χ can naturally be introduced in this framework through the Lagrangian⁴:

$$\begin{aligned} \mathcal{L}_\chi &= \frac{1}{2} i \bar{\chi} \gamma^\mu \partial_\mu \chi - g_\chi \Phi \bar{\chi} \chi + \text{h.c.} \\ &= \frac{1}{2} i \bar{\chi} \gamma^\mu \partial_\mu \chi - \frac{g_\chi}{\sqrt{2}} s \bar{\chi} \chi - i \frac{g_\chi}{\sqrt{2}} a \bar{\chi} \gamma^5 \chi. \end{aligned} \quad (2.4)$$

Once Φ acquires a vacuum expectation value (vev) through the spontaneous symmetry breaking mechanism, $\Phi = \frac{1}{\sqrt{2}}(v_\Phi + s + ia)$ with $v_\Phi^2 = \frac{\mu_\Phi^2}{\lambda_\Phi} + \frac{\epsilon_\Phi^2}{\lambda_\Phi} \simeq \frac{\mu_\Phi^2}{\lambda_\Phi}$, it generates automatically

⁴We will work in the framework of Majorana dark matter throughout this analysis. The extension to the Dirac case is straightforward.

the mass of s ($m_s = \sqrt{2\lambda_\Phi}v_\Phi$), the saa coupling ($\lambda_{saa} = \lambda_\Phi v_\Phi$) and the dark matter mass ($m_\chi = \sqrt{2}g_\chi v_\Phi$). Moreover, the parameter Λ in eq. (2.2) can always be freely chosen up to a redefinition of the couplings c_{ii} 's. Choosing logically $\Lambda = v_\Phi$, being the scale of new physics, one can eliminate v_Φ and re-express the relevant part of the Lagrangian (see eq. (2.2)) as the function of masses and couplings as:

$$\begin{aligned}
-\mathcal{L} \supset & \frac{\sqrt{\lambda_\Phi}C_{GG}}{m_s}sG_{\mu\nu}^\alpha G_{\alpha}^{\mu\nu} + \frac{\sqrt{\lambda_\Phi}C_{GG}}{m_s}aG_{\mu\nu}^\alpha \tilde{G}_{\alpha}^{\mu\nu} \\
& + \frac{\sqrt{\lambda_\Phi}C_{BB}c_W^2}{m_s}sF_{\mu\nu}F^{\mu\nu} + \frac{\sqrt{\lambda_\Phi}C_{BB}c_W^2}{m_s}aF_{\mu\nu}\tilde{F}^{\mu\nu} \\
& + \frac{m_s^2}{2}s^2 + \frac{m_a^2}{2}a^2 + \sqrt{\frac{\lambda_\Phi}{2}}m_s s a^2 + \sqrt{\frac{\lambda_\Phi}{2}}m_s s^3 + \frac{\lambda_\Phi}{4}(s^2 + a^2)^2 \\
& + \frac{1}{2}m_\chi\bar{\chi}\chi + \sqrt{\frac{\lambda_\Phi}{2}}\frac{m_\chi}{m_s}(s\bar{\chi}\chi + ia\bar{\chi}\gamma^5\chi),
\end{aligned} \tag{2.5}$$

where $C_{ii} = 2c_{ii}$, $c_W^2 = \cos^2\theta_W$, $F_{\mu\nu}(\tilde{F}_{\mu\nu})$ is the field(dual field) strength for the photon field, $m_a = \sqrt{2}\epsilon_\Phi$ is the pseudoscalar mass and we have worked with⁵ $c_{WW} \rightarrow 0$. If now one fixes the mass of the scalar at 750 GeV, all the physics is then completely determined by a set of five parameters (C_{GG} , C_{BB} , λ_Φ , m_a , m_χ). If one further wants to fit simultaneously, the width of s , the di-photon cross-section and the relic abundance, one is left with only two free parameters. The model then becomes very predictive.

Notice that, sticking to an effective field theory approach, the dimension-5 interaction terms in eq. (2.2), between the field Φ and the SM gauge bosons, do not preserve the $U(1)$ symmetry. As already mentioned, an explicit breaking of the $U(1)$ symmetry has already been introduced in eq. (2.3) in order to have a non-zero mass for the pseudoscalar field a and thus, the presence of these dimension-5 operators is rather natural. Moreover, in the studied scenario, as will be clarified subsequently, the coefficients of these operators, e.g., c_{BB} , c_{GG} , required to be much suppressed in order to generate an experimentally viable production cross-section with collimated photons. In this setup, one might also consider the inclusion of dimension-6 operators, e.g., $\frac{\tilde{c}_{GG}}{\Lambda^2}\Phi^\dagger\Phi G_a^{\mu\nu}G_{\mu\nu}^a$, which preserve the $U(1)$ symmetry. These terms, with the choice of $\Lambda = v_\Phi$, after spontaneous breaking of the $U(1)$ symmetry would produce new effective dimension-5 interactions for the scalar, featuring the same $1/v_\Phi$ suppression associated with the $U(1)$ violating ones. This would modify the trilinear couplings between the scalar component of Φ and the SM gauge bosons and thus, will affect the production vertex and decay width of the scalar field. Now, as will be explained later, in the studied framework the decay width is dominated by the tree-level saa coupling and hence, the impact of including dimension-6 terms would hardly affect our findings. The relevant interactions for the pseudoscalar, however, remain the same since only dimension-5 operators can lead to couplings linear in the a field.

The main purpose of this paper is to provide a phenomenological explanation of the di-photon excess through the production of collimated photon pairs simultaneously with the main aspects of dark-matter physics. For this reason in our analysis we will freely vary the couplings of various operators and hence, the effect of dimension-6 operators can be straightforwardly encapsulated through a redefinition, including the contributions from different di-

⁵The coupling C_{BB} , from the requirement of gauge invariance, also generates interactions like sZZ , $sZ\gamma$, etc. We, however, do not explore them in detail since they are observed to have a marginal impact in our analysis.

mensionalities, of the effective couplings between the scalar field s and the SM gauge bosons. Thus, we will not explicitly refer any further to the presence of dimension-6 operators.

The existence of dimension-six interactions, nevertheless, would appear effective for the pair production of ss or aa which would lead to characteristic signals, e.g., $gg \rightarrow ss \rightarrow 4a \rightarrow 8\gamma$. A detail and subsequent discussion of such signal, however, is beyond the scopes of the current article.

We note in passing that for this work we consider the width of the resonance in the span of $4 - 60$ GeV, the di-photon production cross-section in the range of $1 - 10$ fb and finally, the relic density $\Omega h^2 \approx 0.12$ [189].

3 The LHC analysis

3.1 Condition to mimic a di-photon process

The decay of a substantially light pseudoscalar can produce highly collimated photons (dubbed as “photon jets”) which can be potentially misidentified as a single photon. The minimal condition to realize this kind of scenario is that the opening angle $\Delta\phi \sim 2/\gamma$ ($\gamma = m_s/2m_a$, being the boost factor) of the two photons emitted in the $a \rightarrow \gamma\gamma$ process remains below the angular resolution of the LHC detectors, which is $\sim \mathcal{O}(20 \text{ mrad})^6$ [190, 191]. This condition provides an upper bound as $m_a \lesssim 2 \text{ GeV}$ [179, 192, 193]. At the same time, a very high boost typically generates an enhanced decay length for a . In order to mimic the LHC di-photon signal, the decay should occur before the electromagnetic calorimeter (ECAL), which is at a distance of ~ 1 meter from the collision point. A minimal requirement is then set on the decay length of a , $l = \beta\gamma/\Gamma_a \ll 1$ meter, where $\beta\gamma = \sqrt{\gamma^2 - 1}$ and Γ_a represents the total decay width for a . This constraint is mostly relevant for values of $m_a \lesssim 3m_{\pi^0}$ (m_{π^0} denotes the neutral pion mass), such that only the decay process $a \rightarrow \gamma\gamma$ remains actually accessible⁷. As a consequence, the decay length, for $m_a \lesssim 3m_{\pi^0}$, can be expressed only as the function of m_a and C_{BB} as:

$$l \approx 7.3 \text{ cm} \left(\frac{m_s}{750 \text{ GeV}} \right)^3 \left(\frac{500 \text{ MeV}}{m_a} \right)^4 \left(\frac{0.05}{\Gamma_s/m_s} \right) \left(\frac{0.01}{C_{BB}} \right)^2. \quad (3.1)$$

Now the requirement of $l \lesssim 1$ meter gives the following lower bound on C_{BB} :

$$C_{BB} \gtrsim 2.7 \times 10^{-3} \left(\frac{m_s}{750 \text{ GeV}} \right)^{3/2} \left(\frac{500 \text{ MeV}}{m_a} \right)^2 \left(\frac{0.05}{\Gamma_s/m_s} \right)^{1/2}. \quad (3.2)$$

We remark that eq. (3.2) should be regarded as a conservative bound. As argued in ref. [181], for example, displaced photons from $a \rightarrow \gamma\gamma$ process corresponding to $l \sim 1 - 10$ cm, might already be distinguishable from a pair of prompt photons originating from $s \rightarrow \gamma\gamma$ process.

We present in figure. 1, for illustration, the allowed m_a values according to the one of C_{BB} , demanding a pair of collimated photons from $a \rightarrow \gamma\gamma$ process. This plot is made with a fixed value of $C_{GG} = 0.0005$ and of λ_Φ , corresponding to $\Gamma_s = 40$ GeV. A further constraint on the associated decay length $l \lesssim 1$ meter disfavors the gray coloured region corresponding to $m_a \lesssim 200$ MeV. For reference we have also plotted the line corresponding to $l = 10$ cm in the (m_a, C_{BB}) plane. We also show the $m_a \lesssim 100$ MeV region, as reported in ref. [194],

⁶This corresponds to $\sim 1.15^\circ$.

⁷One should note that in this regime the invisible branching fraction from $a \rightarrow \bar{\chi}\chi$ process remains negligible, given the suppression of the $a\bar{\chi}\chi$ coupling (see eq. (2.5)) for $m_\chi \lesssim m_a/2$.

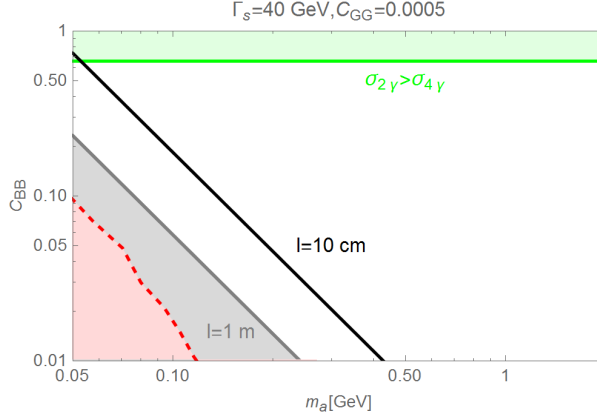


Figure 1. Allowed region in the (m_a, C_{BB}) plane when $a \rightarrow \gamma\gamma$ process leads to a collimated photon pairs and simultaneously, the associated decay length remains $\lesssim 1$ meter (excluding the gray coloured region). The black coloured solid line represents a typical displaced decay with $l = 10$ cm. The red coloured region is excluded by constraints reported in ref. [194]. The green coloured region corresponds to $\sigma_{2\gamma} > \sigma_{4\gamma}$ which is duly explained in the text.

(represented by the red coloured region) which remains constrained from the direct searches of a light pseudoscalar in beam-dump and fixed target experiments.

It is important to note that the process $pp \rightarrow s \rightarrow \gamma\gamma$ is also possible in our model since s , just like a , can couple to two photons through the same coupling C_{BB} . We have thus represented (in green colour) in figure 1 the region where⁸ $\sigma_{2\gamma} \equiv \sigma(pp \rightarrow s \rightarrow \gamma\gamma) > \sigma_{4\gamma} \equiv \sigma(pp \rightarrow s \rightarrow aa \rightarrow 4\gamma)$.

The allowed range, $0.2 \text{ GeV} \leq m_a \leq 2 \text{ GeV}$, can be constrained further by considering suitable isolation cuts. Indeed, experimental searches of the somewhat similar $h \rightarrow aa \rightarrow 4\gamma$ process, with h being the SM Higgs boson, can discriminate, through suitably defined calorimeter variables, isolated photons from the ones coming from π^0 decays. Given the similarity between the $\pi^0 \rightarrow \gamma\gamma$ and $a \rightarrow \gamma\gamma$ processes, it is useful to adopt those variables to reduce the probability of faking single photon signals [200].

We have then redetermined, in figure 2 the allowed parameter space in the (m_a, C_{BB}) plane, for $C_{GG} = 0.001$ and for four values of λ_Φ , ranging from 0.25 to 4 giving Γ_s in the span of 4 GeV to 60 GeV, by using the treatment illustrated in ref. [182]. The upper limit on m_a , compared to figure 1, is now lowered to 500 – 600 MeV (based on the similar considerations a slightly weaker limit of 800 MeV has been found in ref. [181]). Interestingly, most of the allowed parameter space now lies below the threshold of $3m_{\pi^0}$, above which the hadronic decays become accessible for the pseudoscalar. As a consequence, in this regime, $Br(a \rightarrow \gamma\gamma) = 1$ irrespective of the values of C_{BB} and C_{GG} . As will be clarified in the next section, this will lead to a rather predictive scenario. To utilize this predictive behaviour, we will use this region of the parameter space for most of the analytical estimates presented in this work. In the absence of a dedicated experimental study, we have nevertheless included the region $m_{3\pi^0} \leq m_a \leq 2 \text{ GeV}$ in the numerical computations.

We finally remark that the allowed parameter space can be constrained even further by considering the different photon conversion rate, i.e., the probability of the interaction of a

⁸In reality in figure 1 the region corresponding to $C_{BB} \gtrsim 0.1$ is excluded by the 8 TeV searches of $s \rightarrow ZZ, Z\gamma$ processes [195–199]. One should interpret figure 1 just as an illustrative plot.

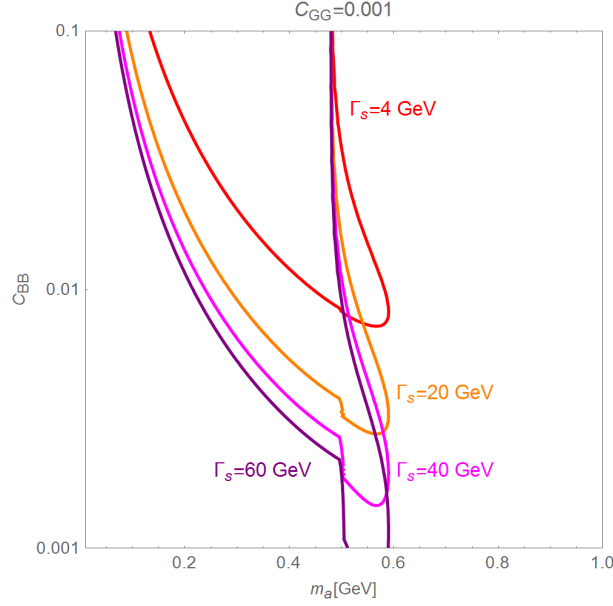


Figure 2. The variations of coupling C_{BB} with m_a for the different values of Γ_s , ensuring that the pseudoscalar decays before the ECAL and produces sufficiently collimated photon pairs to mimic the observed di-photon signal. For this plot $C_{GG} = 0.001$.

photon within the calorimeter to produce a e^+e^- pair, between 2γ and 4γ events. A detailed investigation of this possibility has been performed in ref. [179].

3.2 Di-jet constraint

As already emphasized, that the couplings C_{GG} and C_{BB} are constrained, by the requirement of the relative absence of additional signals, e.g., di-jets, $Z\gamma$ or ZZ etc., compared to the di-photon one. In our scenario the strongest constraint on C_{GG} comes from the di-jet searches. For this purpose we have imposed that $\sigma(pp \rightarrow s \rightarrow jj) < 2.5$ pb, at 8 TeV centre-of-mass⁹ energy [170, 171].

In the $Br(a \rightarrow \gamma\gamma) = 1$ regime the corresponding cross-section is given by:

$$\begin{aligned} \sigma(pp \rightarrow s \rightarrow jj) &= \frac{\pi^2}{8m_s s} I_{GG,8TeV} \frac{2C_{GG}^2 \lambda_\Phi m_s}{\pi} Br(s \rightarrow gg) \\ &\approx \frac{\pi^2}{8m_s s} I_{GG,8TeV} \frac{2C_{GG}^2 \lambda_\Phi m_s}{\pi} \frac{32C_{GG}^2}{1 + 32C_{GG}^2}, \end{aligned} \quad (3.3)$$

where we have used eq. (2.5) and retained the dependence only on the aa and gg channels (see next section) in the determination of Γ_s . The quantities s and $I_{GG,8TeV}$ represent the square of the centre-of-mass energy and the integral of a function involving dimensionless quantities like parton distribution functions (PDFs), evaluated at $\sqrt{s} = 8$ TeV, respectively. It can be easily seen from eq. (3.3) that the condition $\sigma(pp \rightarrow s \rightarrow jj) < 2.5$ pb, with¹⁰ $I_{GG,8TeV} = 280$, is satisfied for $C_{GG} \lesssim 0.1$.

⁹One can, for example, see ref. [201, 202] for a similar bound with 13 TeV centre-of-mass energy.

¹⁰Here we have considered MSTW 2008 NLO PDF [203–205] and included the next-to-leading order scaling factor.

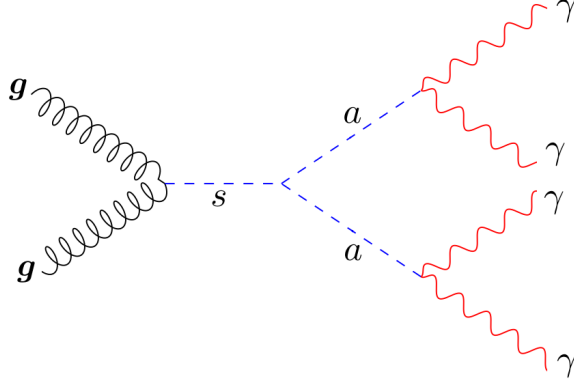


Figure 3. Main channel contributing to the di-photon excess observed at the LHC.

3.3 Fitting the di-photon excess

The production cross-section for the studied di-photon signal through the gluon fusion process $gg \rightarrow s \rightarrow aa \rightarrow 4\gamma$ (see figure 3) is given by

$$\sigma_{4\gamma} = \frac{\pi^2}{8m_s s} \Gamma(s \rightarrow gg) Br(s \rightarrow aa) [Br(a \rightarrow \gamma\gamma)]^2 I_{GG}, \quad (3.4)$$

with I_{GG} being the integral over the PDFs. This is estimated to be $\simeq 3400$ using MSTW 2008 NLO PDF [203–205] and including the next-to-leading order scaling factor in the definition of I_{GG} . The desired signal, as already mentioned, can also be induced through the photon fusion [206] process $\gamma\gamma \rightarrow s \rightarrow aa \rightarrow 4\gamma$. We, however, do not consider this possibility in this article.

At this point one can use eq. (2.5) to compute:

$$\Gamma(s \rightarrow aa) = \frac{\lambda_\Phi m_s}{16\pi}, \quad \Gamma(s \rightarrow gg) = \frac{2\lambda_\Phi m_s}{\pi} C_{GG}^2. \quad (3.5)$$

Now the processes $s \rightarrow \gamma\gamma$, $s \rightarrow ZZ$, $s \rightarrow Z\gamma$ are suppressed compared to $s \rightarrow gg$ mode by factors like $8/\cos^4 \theta_W$, $8/\sin^4 \theta_W$, $8/\sin^2 \theta_{2W}$, respectively, even when $C_{BB} \sim C_{GG}$. The upper limit of C_{BB} is constrained from the photon fusion process. The invisible decay mode $s \rightarrow \bar{\chi}\chi$ also suffers suppression, either from the $s\bar{\chi}\chi$ coupling $\propto m_\chi/m_s$ (for $m_\chi \ll m_s$) or from the phase space factor (when $m_\chi \sim m_s/2$). Thus, collectively one can use eq. (3.5) to get

$$Br(s \rightarrow aa) \approx \frac{1}{1 + 32C_{GG}^2}. \quad (3.6)$$

This equation is especially valid for $m_a \lesssim 500$ MeV, i.e., below the $3m_{\pi^0}$ threshold when $Br(a \rightarrow \gamma\gamma) = 1$. In this region, combining eq. (3.6) with eq. (3.4) one obtains:

$$\sigma_{4\gamma} \approx \frac{4\pi^2}{s} I_{GG} \frac{C_{GG}^2}{(1 + 32C_{GG}^2)^2} \frac{\Gamma_s}{m_s} \simeq 16 \text{ fb} \frac{(\Gamma_s/m_s)}{0.05} \frac{\left(\frac{C_{GG}}{0.001}\right)^2}{(1 + 32C_{GG}^2)^2}. \quad (3.7)$$

Interestingly, for $m_a \lesssim m_{3\pi^0}$ the 4γ production cross-section depends only on C_{GG} and λ_Φ couplings while remains independent of C_{BB} , the $a\gamma\gamma$ coupling. Indeed, from eq. (3.7) one can see the possibility to obtain $\sigma_{4\gamma} \sim \mathcal{O}(1 - 10 \text{ fb})$ together with a large Γ_s ($\sim 5\% m_s$) for

low values of C_{GG} , consistent with the di-jet constraints. This comes trivially from the fact that a large λ_Φ coupling is responsible for the large Γ_s and a sizable $\sigma_{4\gamma}$ without requiring high values of C_{GG} and C_{BB} couplings. We finally notice that the di-jet limit, $C_{GG} \lesssim 0.1$ implies that for all practical purposes $\Gamma_s \approx \Gamma(s \rightarrow aa)$. This feature allows us to trade, in the analytical expressions considered in this section, the free parameter λ_Φ with the physical observable Γ_s .

It is important to note that although the coupling C_{BB} does not directly appear in $\sigma_{4\gamma}$, it is not unconstrained in nature. In fact, the absence of di-photon signal from photon fusion process and the requirement of a pseudoscalar decaying before the ECAL provide an upper and lower bound for C_{BB} , respectively. In reality, however, there exists one more way to constrain the coupling C_{BB} by demanding a suppressed $s \rightarrow \gamma\gamma$ process. We have already used this to derive eq. (3.6). The condition $\frac{\Gamma(s \rightarrow \gamma\gamma)}{\Gamma(s \rightarrow aa)} \ll 1$, implies $C_{BB} \ll 0.65$.

A good fit for $\sigma_{4\gamma}$ can also be obtained for $m_a \gtrsim 500$ MeV with a sizable branching fraction of a into the hadronic channels. In such a case the cross-section is obtained by multiplying eq. (3.7) with $(Br(a \rightarrow \gamma\gamma))^2 \approx (C_{BB}^4 c_W^8 / 64 C_{GG}^4)$ and is given by:

$$\sigma_{4\gamma} \approx 5 \text{ fb} \frac{(\Gamma_s/m_s)}{0.05} \frac{1}{(1 + 32 C_{GG}^2)^2} \left(\frac{C_{BB}}{0.005} \right)^4 \left(\frac{0.005}{C_{GG}} \right)^2. \quad (3.8)$$

The condition for the dominance of $s \rightarrow 4\gamma$ channel compared to $s \rightarrow \gamma\gamma$ mode is similarly modified and becomes:

$$C_{GG} \lesssim 0.27 C_{BB}^{1/2}. \quad (3.9)$$

4 The analysis implementing the dark sector

Once one couples the dark matter to the field Φ through eq. (2.4), one can compute the relic abundance of χ as a function of C_{GG} and λ_Φ for different m_χ values and check if there exists parameter space compatible with the observed di-photon excess. The result is presented in figure 4 where we have shown in the (λ_Φ, C_{GG}) plane the parameter space allowed by the three constraints, namely: (i) $1 \text{ fb} \leq \sigma_{4\gamma} \leq 10 \text{ fb}$, (ii) $4 \text{ GeV} \leq \Gamma_s \leq 60 \text{ GeV}$ and (iii) correct relic abundance ≈ 0.12 [189]. It is evident from figure 4 that for any value of the dark matter mass between 200 GeV to 700 GeV, there exist values of C_{GG} that give a $\sigma_{4\gamma} \sim \mathcal{O}(1 - 10 \text{ fb})$ together with $4 \text{ GeV} \lesssim \Gamma_s \lesssim 60 \text{ GeV}$ while still respecting the PLANCK constraints. We will study the different regimes in detail in the following sub-sections.

4.1 $m_\chi \lesssim 375 \text{ GeV}$

For $m_a < m_\chi \lesssim m_s/2$ the dark matter can annihilate into $\gamma\gamma$, gg and aa , as well as into $Z\gamma$ and ZZ , whether kinematically open (the last two modes are suppressed in the setup considered in this work).

The $\gamma\gamma$ and gg channels mostly originate from s-channel exchange of the pseudoscalar. The corresponding thermally averaged cross-sections, using eq. (3.6) and eq. (3.7), can be

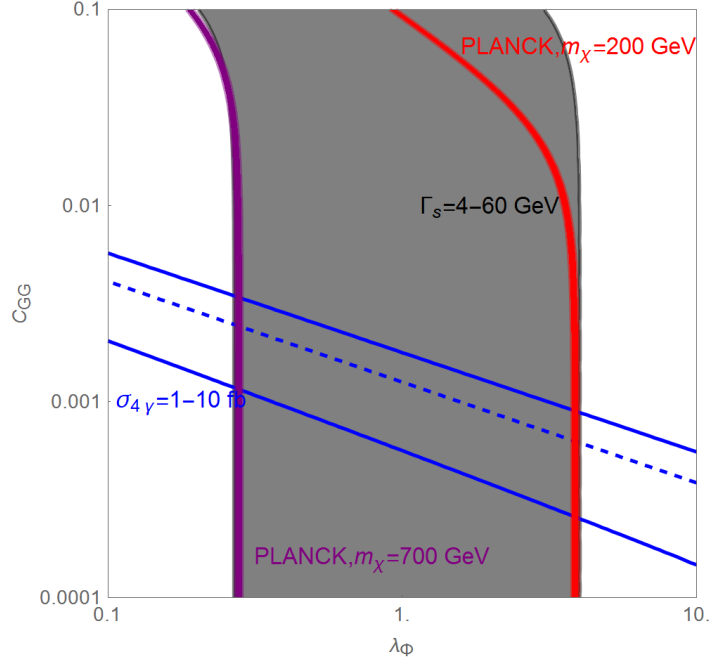


Figure 4. The allowed parameter space in the (λ_Φ, C_{GG}) plane consistent with the observations of: (a) correct relic density, (b) the observed di-photon production cross-section and (c) a relatively large width of the resonance. The gray coloured band corresponds to the large width region (4 – 60 GeV) whereas the blue coloured lines (solid and dashed) represent a di-photon production cross-section of 1 – 10 fb, as observed at the LHC. The red and magenta coloured lines correspond to points that respect the PLANCK constraints for $m_\chi = 200$ GeV and 700 GeV, respectively. For this plot we consider $C_{BB} = 0.005$.

estimated as [47, 173, 174]:

$$\begin{aligned}
 \langle \sigma v \rangle_{gg} &= \frac{16\lambda_\Phi^2 C_{GG}^2 m_\chi^2}{\pi m_s^4} \approx 5 \times 10^{-30} \left(\frac{\sigma_{4\gamma}}{16 \text{ fb}} \right) \left(\frac{\Gamma_s/m_s}{0.05} \right) \left(\frac{m_\chi}{100 \text{ GeV}} \right)^2 \left(\frac{750 \text{ GeV}}{m_s} \right)^4, \\
 \langle \sigma v \rangle_{\gamma\gamma} &= \frac{2\lambda_\Phi^2 C_{BB}^2 c_W^4 m_\chi^2}{\pi m_s^4} \\
 &\approx \frac{9 \times 10^{-27}}{(1 + 32C_{GG}^2)^2} \left(\frac{C_{BB}}{0.1} \right)^2 \left(\frac{\Gamma_s/m_s}{0.05} \right)^2 \left(\frac{m_\chi}{100 \text{ GeV}} \right)^2 \left(\frac{750 \text{ GeV}}{m_s} \right)^4, \quad (4.1)
 \end{aligned}$$

in the unit of cm^3s^{-1} . We now clearly see the difficulty to achieve the correct relic abundance through the s-wave channel due to the constraints on $\sigma_{4\gamma}$ and $C_{BB} \lesssim 0.1$ for the gg and $\gamma\gamma$ final states, respectively. Moreover, the gluon channel is also constrained by the di-jet constraint, eq. (3.3), and, to a lesser extent and only for $m_a \geq 500$ MeV, by the FERMI searches of gamma-rays from DSph [207, 208]. On the other hand, the contribution from the photon channel is, instead, substantially excluded by the FERMI searches of monochromatic gamma-ray signals [209], which give a limit as strong as $10^{-(29 \div 30)} \text{ cm}^3\text{s}^{-1}$ for $m_\chi \lesssim 100$ GeV.

Even if the aa channel gives rise to a p-wave velocity suppressed cross-section (by CP arguments), it is the dominant annihilation process at the decoupling time. It proceeds mostly through the s -channel exchange of the scalar s as depicted in the figure 5. The concerned

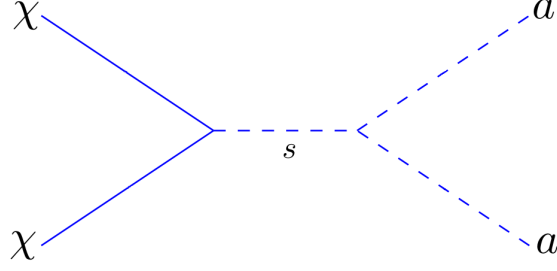


Figure 5. Main channel contributing to the relic abundance for $m_\chi \lesssim 375$ GeV.

thermally averaged cross-section can be estimated as:

$$\begin{aligned} \langle \sigma v \rangle_{aa} &= \frac{9m_\chi^2 \lambda_\Phi^2}{384\pi m_s^4} v^2 \\ &\approx 5.3 \times 10^{-27} \text{ cm}^3 \text{ s}^{-1} \frac{1}{(1 + 32C_{GG}^2)^2} \left(\frac{\Gamma_s/m_s}{0.05} \right)^2 \left(\frac{m_\chi}{100 \text{ GeV}} \right)^2 \left(\frac{750 \text{ GeV}}{m_s} \right)^4. \end{aligned} \quad (4.2)$$

In figure 6 we show the result of complete numerical analysis, performed precisely by determining the dark matter annihilation cross-sections and its relic density through the package **MicrOMEGAs** [210]. Fixed values for the C_{GG} and C_{BB} couplings, as 0.001 and 0.005, respectively, are used for this figure. Here we have plotted the contour of correct relic density in the (m_χ, Γ_s) plane corresponding to different di-photon production cross-sections. We have also used the same methods and tools for figure 4. We notice from figure 6 that, when $m_\chi \lesssim 375$ GeV, the correct relic density can be achieved, through the annihilation channel of figure 5, for $m_\chi \sim \mathcal{O}(200 - 300 \text{ GeV})$ and for $12 \text{ GeV} \leq \Gamma_s \leq 50 \text{ GeV}$ while still fitting the LHC di-photon data. We also notice the typical pole effect, i.e., when m_χ approaches $m_s/2$, the width needs to be narrow to avoid the under-abundance due to thermal broadening effect [211]. On the other hand, being velocity suppressed, this channel cannot account for any indirect detection signal as it will be discussed in sub-section 4.4.

4.2 $m_\chi > 375$ GeV

When $m_\chi > m_s/2$, the annihilation process $\chi\chi \rightarrow sa$, shown in figure 7, becomes kinematically allowed. This channel will dominate the annihilation process compared to aa (or ss when $m_\chi > m_s$) as it is the only channel which is not velocity suppressed. The relevant thermally averaged cross-section is given by:

$$\langle \sigma v \rangle_{sa} \simeq \frac{\lambda_\Phi^2 m_\chi^2}{8\pi m_s^4} \simeq \frac{3 \times 10^{-25} \text{ cm}^3 \text{ s}^{-1}}{(1 + 32C_{GG}^2)^2} \left(\frac{\Gamma_s/m_s}{0.05} \right)^2 \left(\frac{m_\chi}{400 \text{ GeV}} \right)^2 \left(\frac{750 \text{ GeV}}{m_s} \right)^4. \quad (4.3)$$

As evident from eq. (4.3) that in the case of $\Gamma_s/m_s \sim 5\%$, the cross-section exceeds the thermally favoured value and leads, consequently, to an underabundant dark matter. This is also reflected from figure 6 that the correct relic density for $m_\chi > 375$ GeV is achieved only for $\Gamma_s < 12 \text{ GeV}$. The latter, for the chosen values of C_{GG} , C_{BB} couplings, corresponds to a value of $\sigma_{4\gamma}$ on the lower side.

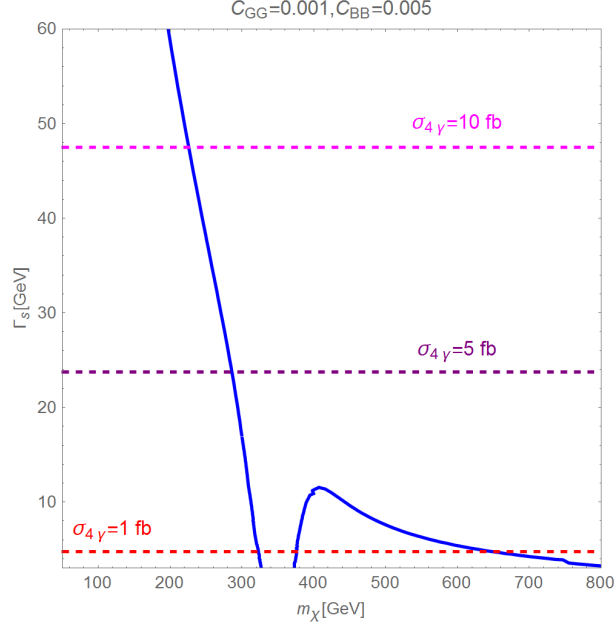


Figure 6. The contour of correct relic density in the (m_χ, Γ_s) plane. The differently coloured horizontal lines correspond to the various values of the di-photon production cross-sections.

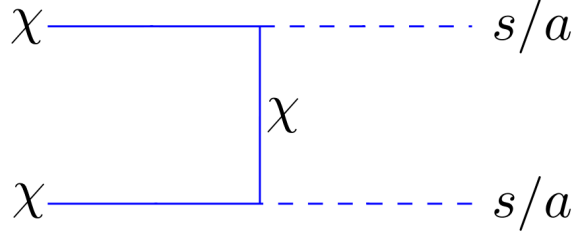


Figure 7. Main channel responsible for the relic abundance for $m_\chi > 375$ GeV.

4.3 Very light case : the freeze in regime

In this case the dominant contribution to the dark matter relic density is given by the annihilations into gg and $\gamma\gamma$ final states through s-channel exchange of the pseudoscalar. The corresponding thermally averaged cross-sections are:

$$\begin{aligned}
\langle\sigma v\rangle_{gg} &= \frac{1024}{\pi} \frac{m_\chi^6}{m_a^4 m_s^4} \lambda_\Phi^2 C_{GG}^2 \\
&\approx 3 \times 10^{-33} \text{ cm}^3 \text{ s}^{-1} \lambda_\Phi^2 C_{GG}^2 \left(\frac{m_\chi}{0.1 \text{ GeV}} \right)^6 \left(\frac{1 \text{ GeV}}{m_a} \right)^4 \left(\frac{750 \text{ GeV}}{m_s} \right)^4, \\
\langle\sigma v\rangle_{\gamma\gamma} &= \frac{128}{\pi} \frac{m_\chi^6}{m_a^4 m_s^4} \lambda_\Phi^2 C_{BB}^2 c_W^4 \\
&\approx 2.2 \times 10^{-34} \text{ cm}^3 \text{ s}^{-1} \lambda_\Phi^2 C_{BB}^2 \left(\frac{m_\chi}{0.1 \text{ GeV}} \right)^6 \left(\frac{1 \text{ GeV}}{m_a} \right)^4 \left(\frac{750 \text{ GeV}}{m_s} \right)^4. \quad (4.4)
\end{aligned}$$

As evident from eq. (4.4) that these cross-sections are very far from the thermally favoured value, $\sim \mathcal{O}(10^{-26}) \text{ cm}^3 \text{ s}^{-1}$, such that the validity of the WIMP paradigm itself

becomes questionable. Thus, we apply the conventional rule-of-thumb of comparing the dark matter annihilation rate $\Gamma_{\text{ann}} = \langle \sigma v \rangle n_{\chi,eq}$ with the Hubble expansion rate, both evaluated at temperature $T \sim \mathcal{O}(m_\chi)$. Considering for reference the $\gamma\gamma$ annihilation channel¹¹, one gets:

$$\frac{\Gamma_{\text{ann}}(T = m_\chi)}{H(T = m_\chi)} \approx \frac{3.2}{(1 + 32C_{GG}^2)^2} \left(\frac{\Gamma_s/m_s}{0.05} \right)^2 \left(\frac{C_{BB}}{0.005} \right)^2 \left(\frac{m_\chi}{0.1 \text{ GeV}} \right)^7 \left(\frac{0.5 \text{ GeV}}{m_a} \right)^4 \left(\frac{750 \text{ GeV}}{m_s} \right)^4. \quad (4.5)$$

For the chosen model parameters with $m_\chi \lesssim 100$ MeV, compatible with the LHC di-photon signal, the dark matter annihilations become inefficient. The dark matter, however, still remains relativistic and the ratio $\frac{\Gamma_{\text{ann}}}{H}$ drops very steeply as long as the dark matter mass is further reduced. The dark matter could nevertheless be kept into thermal equilibrium by the decay process $\Gamma(a \rightarrow \bar{\chi}\chi) = \frac{\lambda_\Phi m_a m_\chi^2}{8\pi m_s^2}$, with a in turn, kept into equilibrium by its couplings with the gauge bosons, at least down to temperatures $\sim \mathcal{O}(m_a)$. This requirement can be translated into the following condition:

$$\begin{aligned} \frac{\Gamma(a \rightarrow \bar{\chi}\chi)(T = m_a)}{H(T = m_a)} &\geq 1, \\ \rightarrow m_\chi &\gtrsim 1.3 \text{ keV} \sqrt{1 + 32C_{GG}^2} \sqrt{\frac{0.02}{\Gamma_s/m_s}} \left(\frac{m_a}{1 \text{ GeV}} \right)^{1/2} \left(\frac{m_s}{750 \text{ GeV}} \right). \end{aligned} \quad (4.6)$$

We thus notice that the $a \leftrightarrow \bar{\chi}\chi$ process is rather efficient and keeps the dark matter in the equilibrium unless $m_\chi \sim \mathcal{O}(1 \text{ keV})$ (or even below) for $\Gamma_s/m_s > 2\%$. Remembering that m_χ below a keV is excluded by structure formation, we conclude that the light dark matter is coupled to the primordial thermal bath until the temperature $T \sim m_a$. After this temperature relativistic decoupling happens with a relic density given by:

$$\Omega_{\chi,\text{rel}} h^2 \approx 9.6 \times 10^{-2} \frac{g_{\text{eff}}}{g_{*,S}(T_d)} \left(\frac{m_\chi}{1 \text{ eV}} \right), \quad (4.7)$$

where $g_{*,S}(T_d)$ represents the number of relativistic degrees of freedom at the temperature T_d of decoupling of the dark matter while g_{eff} is the number of internal degrees of freedom of the dark matter itself. The correct relic density can be achieved for $m_\chi \sim 100$ eV, inconsistently with the conditions stated above. For viable values of the mass, i.e., $\gtrsim \mathcal{O}(\text{keV})$, the dark matter is overabundant. For $\Gamma_s/m_s < 2\%$, on the contrary, there exists a small window of masses, $m_\chi \sim 1 - 10 \text{ keV}$, for which the dark matter cannot get into the thermal equilibrium in the Early Universe. In this case the dark matter is produced through freeze-in [212–214] by the decay of the pseudoscalar. The corresponding relic density is given by:

$$\begin{aligned} \Omega_{FI} h^2 &= \frac{1.09 \times 10^{27} g_a}{g_{*,S}(T = m_a)^{3/2}} \frac{m_\chi \Gamma(a \rightarrow \bar{\chi}\chi)}{m_a^2} \\ &\approx 0.3 \frac{1}{(1 + 32C_{GG}^2)} \left(\frac{\Gamma_s/m_s}{0.01} \right) \left(\frac{m_\chi}{1 \text{ keV}} \right)^3 \left(\frac{750 \text{ GeV}}{m_s} \right)^2 \left(\frac{1 \text{ GeV}}{m_a} \right). \end{aligned} \quad (4.8)$$

If the dark matter is not in the thermal equilibrium, the $s \rightarrow \bar{\chi}\chi$ decay process can also give rise to freeze-in production. However, its contribution is suppressed with respect to

¹¹As m_χ decreases, the relevant annihilation processes occur below the temperature of the QCD phase transition, $\sim \mathcal{O}(200 \text{ MeV})$, and then the hadronic annihilation channel remains inaccessible.

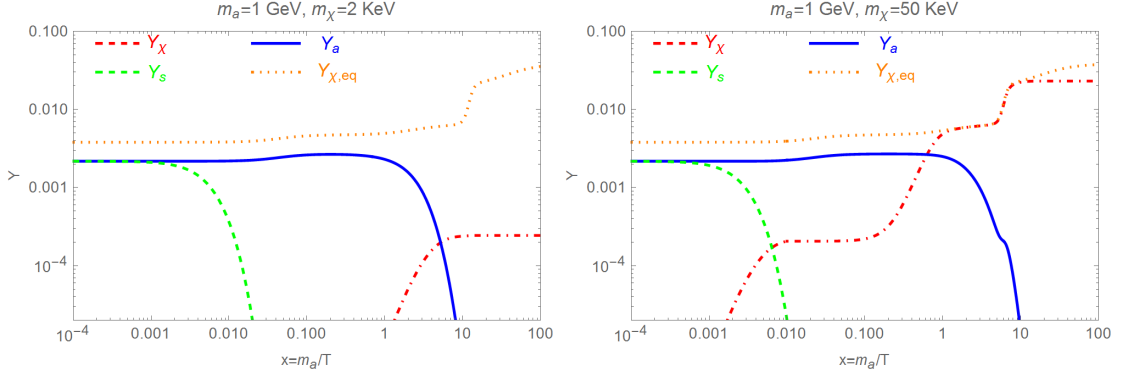


Figure 8. Evolution of the abundance of the s , a and χ fields, as determined by the numerical solution of a system of coupled Boltzmann equations in the very light dark matter regime. The two panels refer to two values of m_χ , namely, 2 keV (left) and 50 keV (right). As evident from the comparison with the dark matter equilibrium distribution (dotted orange coloured line), that for the lighter mass, the dark matter is not capable of getting into the thermal equilibrium and it is produced through freeze-in. In the case of $m_\chi = 50$ keV, the dark matter can however go into the thermal equilibrium and get decoupled while still being relativistic.

eq. (4.8), achieved with the replacement $m_a \rightarrow m_s$. As can be seen, the correct dark matter relic density requires a Γ_s on the smaller side, i.e., $\lesssim 7.5$ GeV giving $\Gamma_s/m_s \lesssim 1\%$ or $m_a \sim 2$ GeV.

The analytical expressions reported above have been validated by solving a system of coupled Boltzmann equations tracking the abundance of the dark matter itself as well as the ones of the scalar field s and of the pseudoscalar field a .

Two examples of the numerical solution of this system are shown in figure 8. Here two values of the dark matter masses have been considered, namely, 2 keV (left) and 50 keV (right). In the lighter dark matter case, we notice that the dark matter abundance is always well below the equilibrium one, represented by the dotted orange coloured¹² curve. It is produced at $T \sim \mathcal{O}(m_a)$, as expected for the freeze-in mechanism, and the numerical value of its abundance, as determined by the solution of the system of equations, is in very good agreement with the theoretical prediction stated in eq. (4.8). We also notice that, for the chosen set of parameters, the dark matter abundance matches the experimental expectation. On the contrary, for the heavier value of the mass, i.e., 50 keV, the dark matter abundance perfectly tracks the thermal equilibrium function until, again, $T \sim \mathcal{O}(m_a)$ and then remains constant confirming the prediction of a relativistic decoupling. Interesting signatures of such a scenario would be the interpretation of the 3.5 keV monochromatic signal as recently observed in different clusters of galaxies [215].

4.4 Dark Matter Detection

One of the main characteristics of the proposed scenario is the fact that here the observed di-photon production cross-section is obtained for relatively lower values of the C_{GG} and C_{BB} couplings, compared to the direct production from a 750 GeV resonance. As already mentioned in the introduction, such values can potentially evade the constraints from the dark matter searches. Regarding direct detection of the dark matter the most relevant interactions

¹²The raise of the curve at high values of x is due to the change of $g_{*,S}$, coming from the QCD phase transition occurring at $T \sim \mathcal{O}(100 \text{ MeV})$. The latter corresponds to $x \sim \mathcal{O}(10)$.

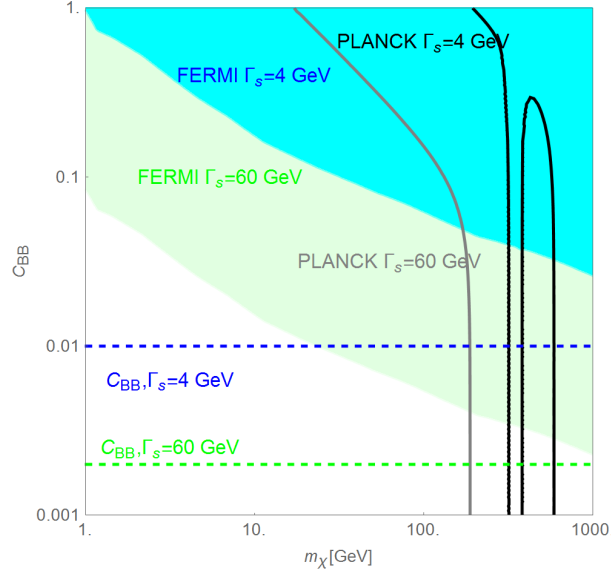


Figure 9. Allowed region in the (m_χ, C_{BB}) plane using the constraints from the FERMI/HESS searches of the gamma-ray lines. The cyan and light green coloured regimes, excluded from the FERMI/HESS searches, correspond to $\Gamma_s = 4$ and 60 GeV, respectively. The black and gray coloured contours represent the correct relic density for the same two values of Γ_s . The blue and green coloured horizontal dashed lines represent two reference C_{BB} values for the same two Γ_s values (see figure 2), which assure a pseudoscalar decay length $\lesssim 1$ meter for $m_a \lesssim 500$ MeV.

are the Spin Independent (SI) ones, produced by the exchange of the scalar field s , as noted also in refs. [47, 173–175]. The SI cross-section, in the studied framework, can be written as:

$$\sigma_{\chi p}^{\text{SI}} \approx 6.7 \times 10^{-48} \text{ cm}^2 \frac{1}{(1 + 32C_{GG}^2)^2} \left(\frac{\Gamma_s/m_s}{0.05} \right)^2 \left(\frac{750 \text{ GeV}}{m_s} \right)^8 \left(\frac{m_\chi}{100 \text{ GeV}} \right)^2 \left(\frac{C_{GG}}{0.001} \right)^2, \quad (4.9)$$

which remains well below the current limit from LUX [216], given the very small value of the C_{GG} coupling. The latter also implies an very suppressed monojet production cross-section and thereby, easily evades the corresponding experimental constraint.

On the contrary, a potential impact is still retained by the indirect dark matter searches. At this moment, the most effective constraint is the one coming from the FERMI and HESS searches of the gamma-ray lines. These searches can trace C_{BB} values $\sim \mathcal{O}(0.001)$ or even smaller. The latter, especially in the low m_a regime, may appear incompatible with the requirement of a pseudoscalar decay length $\lesssim 1$ meter, as depicted in figure 1.

The impact of the indirect detection on the relevant model parameters is represented in figure 9. Here we have shown the excluded regions in the (m_χ, C_{BB}) plane from the FERMI/HESS searches for the minimum and maximum Γ_s values considered in this work, i.e., 4 GeV and 60 GeV, respectively. The regions in the (m_χ, C_{BB}) plane excluded by the FERMI searches, corresponding to $\Gamma_s = 4$ GeV and 60 GeV, are represented with the cyan and light green colour, respectively. The black and the gray coloured contours correspond to the correct relic density, as suggested by the PLANCK. It is apparent from figure 9 that one can extract a *decreasing* minimum value for C_{BB} depending on the *increasing* m_χ value. As a reference, we also compare these limits with the ones obtained from figure 2. The latter are represented by the dashed blue coloured line (for $\Gamma_s = 4$ GeV) and by dashed green coloured

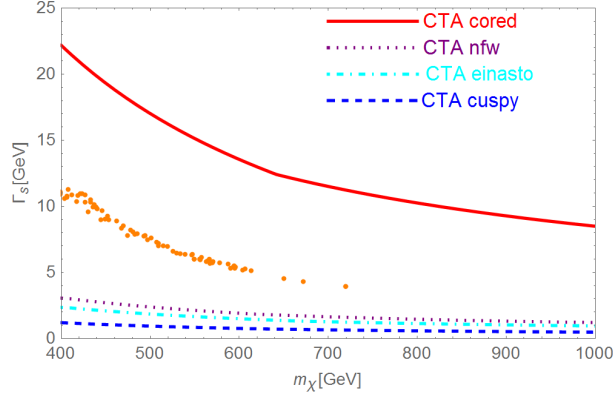


Figure 10. Plot showing the CTA detection prospects in the (m_χ, Γ_s) plane for the gamma-ray boxes originating from the annihilation process $\bar{\chi}\chi \rightarrow sa$. This plot is prepared for the four possible choices of the dark matter profile, as mentioned in the plot. The orange coloured points represent the set of viable model points which emerged from the parameter scan described in the sub-section 4.4.

line (for $\Gamma_s = 60$ GeV), respectively. One should note that the C_{BB} values considered for figure 9 always give a pseudoscalar decay length $\lesssim 1$ meter as observed in figure 2.

Concerning detection of the dark matter, an interesting scenario appears for $m_\chi > m_s/2$. In this case the s-wave annihilation process $\bar{\chi}\chi \rightarrow sa$ gives a characteristic signal represented by a wide gamma-ray box [217]¹³ which can be probed by the CTA [218] in the near future. The detection prospects are shown in figure 10. It is evident from figure 10 that the CTA can efficiently probe the studied corner of the parameter space and thus, in the absence of any signal, can also exclude the entire viable parameter space.

4.5 Summary

We summarized our results in figure 11. This plot accommodates a set of points in the (m_χ, Γ_s) plane possessing the correct relic density and corresponds to different values of the di-photon production cross-sections, as represented with the three different colours. This result is obtained after a dedicated numerical scan on the set of parameters $(C_{GG}, C_{BB}, \lambda_\Phi, m_\chi, m_a)$ over the following ranges, keeping m_s fixed at 750 GeV:

$$\begin{aligned} C_{GG} &\in [5 \times 10^{-4}, 0.1], \quad C_{BB} \in [5 \times 10^{-4}, 0.1], \\ m_\chi &\in [10, 1000] \text{ GeV}, \quad m_a \in [0.2, 2] \text{ GeV}, \\ \lambda_\Phi &\in [0.25, 4]. \end{aligned} \quad (4.10)$$

Finally, we kept only the points respecting

$$1 \text{ fb} \lesssim \sigma_{4\gamma} \lesssim 10 \text{ fb}, \quad 4 \text{ GeV} \lesssim \Gamma_s \lesssim 60 \text{ GeV}, \quad \Omega h^2 \approx 0.12, \quad (4.11)$$

and remain compatible with the relevant accelerator constraints (8 TeV searches of the di-jets, $Z\gamma$, ZZ , decay length of the pseudoscalar $\lesssim 1$ meter etc.) as well as with the dark matter detection limits discussed in the previous sub-section. We can clearly distinguish three regions in figure 11: (i) The low mass region ($m_\chi \lesssim 300$ GeV) in which the relic abundance

¹³A wide gamma-ray box is actually achieved only when $m_\chi > 400$ GeV. The gamma-ray box is narrow when close to the kinematical threshold and becomes sensitive to the gamma-ray line searches. We have included the corresponding constraint in the parameter scan described in the sub-section 4.4.

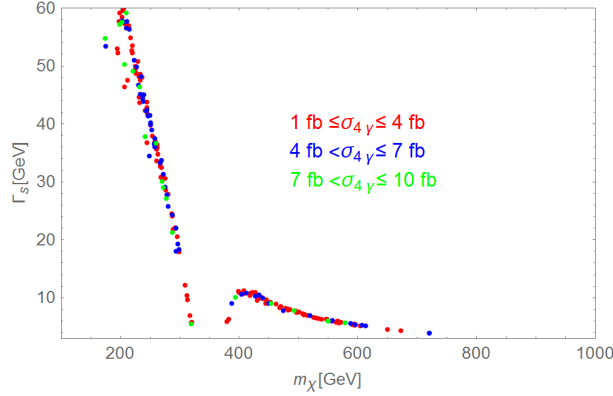


Figure 11. Summary plot representing points in the (m_χ, Γ_s) plane that simultaneously respect the LHC and cosmological constraints and correspond to different di-photon production cross-sections, as depicted with the three different colours.

is dominated by the p-wave (velocity suppressed) annihilation channel $\bar{\chi}\chi \rightarrow s \rightarrow aa$ (see figure 5). In this case a large width ($\Gamma_s \gtrsim 20$ GeV) is necessary to compensate the velocity suppression effect and to avoid *overabundance* of the dark matter in the Universe. (ii) In the second region, $300 \text{ GeV} \lesssim m_\chi \lesssim 400 \text{ GeV}$, the correct abundance is obtained through the pole region of the same diagram (figure 5). In this case, a narrow width of s ($\lesssim 12$ GeV) is necessary to avoid *underabundance* of the dark matter. (iii) Finally, in the last region ($m_\chi \gtrsim 400 \text{ GeV}$), the t-channel (s-wave) annihilation depicted in figure 7 determines the relic abundance of χ . As the final state sa has an odd parity, there is no velocity suppression and a moderate width ($\sim \mathcal{O}(10 \text{ GeV})$) is needed to respect the WMAP/PLANCK constraint.

It is also interesting to notice from figure 11 that if the future ATLAS/CMS analysis determines the value of the width Γ_s more precisely, it will become possible to estimate the dark matter mass, consistent with the cosmological constraint. If the large width scenario is confirmed, an electroweak scale mass will be favoured.

Before moving to the conclusion we just briefly comment on some theoretical aspects relative to our scenario. Our work is motivated by pure phenomenological purposes and thus, we have not referred to any particular setup for the origin of the couplings C_{GG} and C_{BB} . As pointed out, for example, in refs. [20, 25, 182, 219, 220], the extra matter needed at the loop level to generate the C_{GG}, C_{BB} couplings affects the running of the gauge couplings, driving them towards the non-perturbative regime already at energy scales of a few TeV. Thus, too much high values of these couplings remain disfavoured. Similar to ref. [182], we then remark that the low values of the C_{GG}, C_{BB} couplings needed to fit the observed LHC excess, compared to the case of a resonance decaying into two isolated photons, relax also the theoretical constraints, besides the experimental ones. In scenarios like the one considered in this work, a strong impact is still retained by the renormalization group equation running of the quartic coupling λ_Φ . In view of this, the higher values of λ_Φ considered here, corresponding to $\Gamma_s \sim 40 - 60 \text{ GeV}$, might result in tension with constraints from the theoretical consistency. On the other hand, we remark that except for $m_\chi \lesssim 300 \text{ GeV}$, the phenomenological constraints favour values of Γ_s below 12 GeV. The latter corresponds to $\lambda_\Phi \lesssim 0.8$ which allows to maintain a safe theoretical framework at least up to an energy scale relevant for the phenomenological analysis.

5 Conclusion

We have considered in this note the possibility that the di-photon signal is in reality produced through the decay of two light pseudoscalars into two pairs of highly collimated photons, indistinguishable from the isolated photons. We have shown that, in a spontaneous symmetry breaking framework, once one couples the scalar sector to the dark matter, it appears feasible to re-open a large range of the dark matter masses compatible with all the present experimental constraints.

Acknowledgments

Y.M. wants to thank especially J.B de Vivie whose help was fundamental throughout our work. The authors thanks Emilian Dudas and Ulrich Ellwanger for fruitful discussions. This work is also supported by the Spanish MICINN's Consolider-Ingenio 2010 Programme under grant Multi-Dark **CSD2009-00064**, the contract **FPA2010-17747**, the France-US PICS no. 06482 and the LIA-TCAP of CNRS. Y. M. and G. A. acknowledges partial support from the European Union FP7 ITN INVISIBLES (Marie Curie Actions, **PITN-GA-2011- 289442**) and the ERC advanced grants Higgs@LHC and MassTeV. G. A. thanks the CERN theory division for the hospitality during part of the completion of this project. P.G. acknowledges the support from P2IO Excellence Laboratory (LABEX). This research was also supported in part by the Research Executive Agency (REA) of the European Union under the Grant Agreement **PITN-GA2012-316704** ("HiggsTools").

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